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PHASES OF RANDOM ANTIFERROMAGNETIC SPIN-1 CHAINS

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(today)

Abstract

We formulate a real-space renormalization scheme that allows the study of the effects of bond randomness in the Heisenberg antiferromagnetic spin-1 chain. There are four types of bonds that appear during the renormalization flow. We implement numerically the decimation procedure. We give a detailed study of the probability distributions of all these bonds in the phases that occur when the strength of the disorder is varied. Approximate flow equations are obtained in the weak-disorder regime as well as in the strong disorder case where the physics is that of the random singlet phase.

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I. INTRODUCTION

The effect of quenched impurities on the physics of one-dimensional spin systems is an important and unsolved problem. Many spin chains can be doped chemically and this creates some kind of disorder in the system. In addition the spin-1/2 chain is equivalent to a system of spinless fermions through the Jordan-Wigner transformation. This means that the problem of interacting spinless fermions in a disordered potential is equivalent to a random spin chain problem. There are not so many techniques that allow the study of these systems. The real-space renormalization group is prominent among them. Some time ago a pioneering study by Ma and Dasgupta [1] showed that the spin-1/2 Heisenberg antiferromagnetic chain with bond randomness is in a so-called random-singlet phase. In this phase the spins are locked into singlets that extend over arbitrarily long distances, in a pattern dictated by the bond distribution. It has been realized recently that the results of their renormalization procedure are in fact *exact* [2]. This random-singlet phase may capture the physics of higher-dimensional disordered systems [3].

In the spin-1/2 case, the random-singlet phase appears for various kind of disorder and in a wide regions of the phase diagram when one adds XXZ anisotropy. This results from the study of the weak disorder regime by bosonizing the spin chain [4].

The spin-1 Heisenberg antiferromagnetic chain has a physics which is vastly different in the pure case. There is a gap to spin excitations and a finite spin-spin correlation length. These features can be best understood by consideration of a hidden topological order [5,6]. In fact the ground state of the spin-1 chain has a hidden long-range order that can be measured only by use of a nonlocal correlation function, the so-called string order parameter. It is a natural question to ask what happens to these peculiar features under the influence of disorder. In fact the original Ma-Dasgupta renormalization scheme requires a broad enough bond distribution to work [7]. So more complex schemes have been proposed [8–10]. As a function of the disorder strength, it has been established that there is a phase transition between a low-disorder gapless phase with hidden order and a strong disorder phase which is the random-singlet phase of Ma and Dasgupta (gapless and no hidden order).

In this paper, we give a detailed construction of a renormalization scheme suited to the study of the spin-1 chain. We generalize the Ma-Dasgupta decimation procedure by keeping more degrees of freedom. A brief account has been given in Letter form [10]. Here we obtain explicit flow equations that are valid deep inside each of the phases that appear. We are able to follow the spin populations as a function of the renormalization scale as well as the evolution of distribution functions of the various kinds of bonds that appear. In section II, we define the renormalization scheme. In section III, we study the weak-disorder phase of the spin-1 chain. Section IV contains our results for the strong disorder regime. The critical regime is studied in section V and section VI contains our conclusions.

II. REAL SPACE RENORMALIZATION PROCEDURE FOR DISORDERED ANTIFERROMAGNETIC SPIN-1 CHAIN

In this section, we explain how to obtain a real space renormalization scheme adequate to study the disordered antiferromagnetic spin-1 chain.

A. The Ma-Dasgupta real-space renormalization in the spin-1/2 case

Ma and Dasgupta have introduced a real-space renormalization procedure for the random antiferromagnetic spin-1/2 chain described by the Hamiltonian

$$H = \sum_i J_i \vec{S}_i \cdot \vec{S}_{i+1} , \quad (2.1)$$

where $\{\vec{S}_i\}$ are quantum spin-1/2 operators and $\{J_i\}$ are positive random variables distributed with some probability distribution $P_0(J)$. Suppose that J_1 is the largest coupling in the chain. The one-bond Hamiltonian,

$$h_0 = J_1 \vec{S}_1 \cdot \vec{S}_2 = \frac{J_1}{2} \left[(\vec{S}_1 + \vec{S}_2)^2 - \vec{S}_1^2 - \vec{S}_2^2 \right] = \frac{J_1}{2} \left[(\vec{S}_1 + \vec{S}_2)^2 - \frac{3}{2} \right] , \quad (2.2)$$

admits two energy levels labeled by $s = 0, 1$

$$e_s = \frac{J_1}{2} \left[s(s+1) - \frac{3}{2} \right] \quad (2.3)$$

the level e_s being $(2s+1)$ times degenerate : $e_0 = -\frac{3}{4}J_1$ represents the singlet, and $e_1 = \frac{1}{4}J_1$ the triplet. At energies much lower than J_1 , the spins \vec{S}_1 and \vec{S}_2 will therefore be frozen into the singlet state $s = 0$. The decimation procedure consists in eliminating the spins \vec{S}_1 and \vec{S}_2 , and in replacing the four spin segment Hamiltonian $H_{0,1,2,3}$ involving the decimated spins \vec{S}_1 and \vec{S}_2

$$H_{0,1,2,3} = h_0 + h_1 \quad \text{where } h_1 = J_0 \vec{S}_0 \cdot \vec{S}_1 + J_2 \vec{S}_2 \cdot \vec{S}_3 \quad (2.4)$$

by the effective Hamiltonian for the remaining spins \vec{S}_0 and \vec{S}_3

$$H_{0,3}^{eff} = E'_{0,3} + J'_0 \vec{S}_0 \cdot \vec{S}_3 \quad (2.5)$$

which is meant to reproduce the four low-energy states of $H_{0,1,2,3}$ which are separated from the other twelve states of $H_{0,1,2,3}$ by a big gap of order J_1 . Using second order perturbation theory to treat h_1 gives

$$E'_{0,3} = -\frac{3}{4}J_1 - \frac{3}{16J_1}(J_0^2 + J_2^2) \quad (2.6)$$

and

$$J'_0 = \frac{J_0 J_2}{2J_1} \quad (2.7)$$

The same procedure may be iterated and successively applied to the new strongest bond of the chain. This defines a flow for the probability distribution of couplings $P(J, \Omega)$ where Ω is the current strongest coupling [1]

$$-\frac{\partial P(J, \Omega)}{\partial \Omega} = P(\Omega, \Omega) \int_0^\Omega dJ_a \int_0^\Omega dJ_b P(J_a, \Omega) P(J_b, \Omega) \delta \left(J - \frac{J_a J_b}{2\Omega} \right) \quad (2.8)$$

This flow equation has to be supplied by some initial condition $P(J, \Omega_0)$. Fisher has shown [2] that, for generic initial conditions, in the reduced variables $\Gamma = \ln\left(\frac{\Omega_0}{\Omega}\right)$ and $z = \frac{1}{\Gamma} \ln\left(\frac{\Omega}{J}\right)$, the probability distribution $R(z, \Gamma)$ of the variable z flows towards the unique fixed point $R^*(z)$

$$R(z, \Gamma) \xrightarrow{\Gamma \rightarrow \infty} R^*(z) \equiv \theta(z) e^{-z} \quad (2.9)$$

where θ is the Heaviside step function. This so-called Random Singlet Fixed Point corresponds to a power-law distribution in the original variables

$$P^*(J, \Omega) = \theta(\Omega - J) \frac{\alpha(\Omega)}{\Omega} \left(\frac{J}{\Omega}\right)^{\alpha(\Omega)-1} \quad \text{where } \alpha(\Omega) \underset{\Omega \ll \Omega_0}{\simeq} \frac{1}{\ln\left(\frac{\Omega_0}{\Omega}\right)} \quad (2.10)$$

for which two typical bonds are typically much weaker than the strongest one Ω . The approximation involved in the use of perturbation theory to obtain the rule (2.7) therefore becomes better and better as the decimation proceeds, and the whole procedure is therefore completely consistent even if the initial distribution is not broad. The Ma-Dasgupta renormalization scheme is moreover very appealing because it gives an interesting physical picture of the random spin-1/2 chain : at low energy, the chain is made of pairs of spins that are coupled together into singlets over arbitrarily long distances, the long singlets bonds being typically much weaker than the smaller ones.

B. Renormalization of an AF bond between two spin-1

The one-bond hamiltonian

$$h_0 = J_1 \vec{S}_1 \cdot \vec{S}_2 = \frac{J_1}{2} \left[(\vec{S}_1 + \vec{S}_2)^2 - \vec{S}_1^2 - \vec{S}_2^2 \right] = \frac{J_1}{2} \left[(\vec{S}_1 + \vec{S}_2)^2 - 4 \right] \quad (2.11)$$

admits three energy levels labeled by $s = 0, 1, 2$

$$e_s = \frac{J_1}{2} [s(s+1) - 4] \quad (2.12)$$

the level e_s being $(2s+1)$ times degenerate : $e_0 = -2J_1$ represents the singlet, $e_1 = -J_1$ the triplet and $e_2 = J_1$ the quintuplet.

In the Ma-Dasgupta procedure, there are only two levels, and “projecting onto the lowest level” is equivalent to “projecting out the highest level”. Here these two possibilities are not equivalent. The first possibility has already been considered in refs. ([8], [7]) where it is shown that the generalization of equation (2.7) describing the effective coupling between \vec{S}_0 and \vec{S}_3 resulting from the projection onto the singlet formed by \vec{S}_1 and \vec{S}_2 , reads

$$J'_0 = \frac{4}{3} \frac{J_0 J_2}{J_1} \quad (2.13)$$

The coefficient $\frac{4}{3}$ being bigger than 1, this rule is not automatically consistent : indeed, the inequalities $J_0 < J_1$ and $J_2 < J_1$ are not sufficient to imply that the new coupling J'_0 is

smaller than the decimated coupling J_1 , in contrast with the rule (2.7) concerning spin-1/2. This procedure can however be considered as qualitatively correct for very broad initial randomness, where the cases which would produce a new coupling J'_0 bigger than the decimated coupling J_1 are statistically negligible. So the strongly disordered antiferromagnetic spin-1 chains are described by the same random singlet fixed point already found in the study of disordered spin-1/2 chains.

For weak initial randomness however, this naive procedure cannot be made consistent. We thus generalize the Ma-Dasgupta procedure with the interpretation of “projecting out the highest level” instead of “projecting onto the lowest level”. More precisely for the antiferromagnetic bond described by the Hamiltonian h_0 , we project out the quintuplet e_2 but to keep the singlet e_0 and the triplet e_1 by replacing the two spin-1 \vec{S}_1 and \vec{S}_2 by two spin-1/2 \vec{S}'_1 and \vec{S}'_2 , and by replacing h_0 by the effective Hamiltonian

$$h_0^{eff} = -\frac{5J_1}{4} + J_1 \vec{S}'_1 \cdot \vec{S}'_2 \quad (2.14)$$

The four spin segment Hamiltonian $H_{0,1,2,3}$ containing the old spins \vec{S}_1 and \vec{S}_2

$$H_{0,1,2,3} = h_0 + h_1 \quad \text{where } h_1 = J_0 \vec{S}_0 \cdot \vec{S}_1 + J_2 \vec{S}_2 \cdot \vec{S}_3 \quad (2.15)$$

has to be replaced by an effective Hamiltonian involving the spins \vec{S}_0 , \vec{S}'_1 , \vec{S}'_2 and \vec{S}_3

$$H_{0,1,2,3}^{eff} = h_0^{eff} + h_1^{eff}. \quad (2.16)$$

If we use a first-order perturbation theory to treat h_1 , we find that the singlet of h_0 remains unchanged, whereas the degeneracy of the triplet is lifted by the perturbation h_1 . Using the Wigner-Eckart theorem for vectorial operators, we find more explicitly that the perturbation h_1 is equivalent to

$$h_1^{eq} = \left(\frac{1}{2} J_0 \vec{S}_0 + \frac{1}{2} J_2 \vec{S}_3 \right) \cdot (\vec{S}_1 + \vec{S}_2) \quad (2.17)$$

We therefore have to choose the effective Hamiltonian

$$h_1^{eff} = J_0 \vec{S}_0 \cdot \vec{S}'_1 + J_2 \vec{S}'_2 \cdot \vec{S}_3 \quad (2.18)$$

since it is equivalent at first-order perturbation theory, using again Wigner-Eckart theorem, to the hamiltonian

$$\left(\frac{1}{2} J_0 \vec{S}_0 + \frac{1}{2} J_2 \vec{S}_3 \right) (\vec{S}'_1 + \vec{S}'_2) \quad (2.19)$$

We have now enlarged the initial space since the chain now contains not only spin-1 but also spin-1/2. However it is possible to define a decimation procedure that is “closed” inside a particular set of spin chains as we will see in the following.

C. The real-space renormalization procedure

We consider the enlarged set of spin chains described by the Hamiltonian

$$H = \sum_i J_i \vec{S}_i \cdot \vec{S}_{i+1} \quad (2.20)$$

where the spin \vec{S}_i is a spin operator of size $s_i = \frac{1}{2}$ or $s_i = 1$, and where the couplings $\{J_i\}$ can be either positive or negative, but have to satisfy the following constraint : for any pair $\{i, j\}$ such that $i < j$, the classical magnetization of the classical ground state of the segment (i, j) , must be smaller or equal to one in absolute value

$$|m_{i,j}| \leq 1 \quad (2.21)$$

where the quantity $m_{i,j}$ reads

$$m_{i,j} = s_i + \sum_{n=i+1}^j s_n \times \text{sign} \left[\prod_{p=i}^{n-1} (-J_p) \right]. \quad (2.22)$$

This condition for $j = i + 1$ gives immediately that there are exactly four types of bonds

- 1) Link of type 1 : Ferromagnetic bond between two spin-1/2
- 2) Link of type 2 : Antiferromagnetic bond between two spin-1/2
- 3) Link of type 3 : Antiferromagnetic bond between one spin-1 and one spin-1/2
- 4) Link of type 4 : Antiferromagnetic bond between two spin-1

Our decimation procedure is the following :

To each bond $(\vec{S}_i, \vec{S}_{i+1}, J_i)$ we associate the energy difference between the higher state and the lower state of the reduced Hamiltonian $J_i \vec{S}_i \cdot \vec{S}_{i+1}$

$$\Delta_i = -J_i \quad \text{if the bond } i \text{ is of type 1} \quad (2.23)$$

$$\Delta_i = J_i \quad \text{if the bond } i \text{ is of type 2} \quad (2.24)$$

$$\Delta_i = \frac{3}{2}J_i \quad \text{if the bond } i \text{ is of type 3} \quad (2.25)$$

$$\Delta_i = 3J_i \quad \text{if the bond } i \text{ is of type 4} \quad (2.26)$$

We pick up the bond $(\vec{S}_{i_1}, \vec{S}_{i_2}, J_{i_1})$ corresponding to the strongest Δ_i of the chain. To define the renormalization rule for this bond, we again divide the four-spin Hamiltonian into

$$H_{i_0, i_1, i_2, i_3} = h_0 + h_1 \quad \text{where } h_0 = J_{i_1} \vec{S}_{i_1} \cdot \vec{S}_{i_2} \quad \text{and} \quad h_1 = J_{i_0} \vec{S}_{i_0} \cdot \vec{S}_{i_1} + J_{i_2} \vec{S}_{i_2} \cdot \vec{S}_{i_3} \quad (2.27)$$

and treat h_1 as a perturbation of h_0 to find the effective Hamiltonian replacing H_{i_0, i_1, i_2, i_3} when the highest energy state of h_0 is removed. We have now to distinguish the four types of bonds

Rule 1) **F bond between two spin-1/2**

The hamiltonian $h_0 = J_{i_1} \vec{S}_{i_1} \vec{S}_{i_2}$ admits two energy levels : the triplet $e_1 = -\frac{|J_{i_1}|}{4}$ and the singlet $e_0 = \frac{3|J_{i_1}|}{4}$. The perturbation h_1 lifts the degeneracy of the triplet, and using Wigner-Eckart theorem, we find that h_1 is equivalent at first order of perturbation theory to

$$h_1^{eq} = \left(\frac{1}{2} J_{i_0} \vec{S}_{i_0} + \frac{1}{2} J_{i_2} \vec{S}_{i_3} \right) \cdot (\vec{S}_{i_1} + \vec{S}_{i_2}) \quad (2.28)$$

To eliminate the singlet state and only keep the triplet state of h_0 , we remove the two spin-1/2 \vec{S}_{i_1} and \vec{S}_{i_2} and replace them by a single spin-1 \vec{S}'_{i_1} , and we replace H_{i_0, i_1, i_2, i_3} by

$$H_{i_0, i'_1, i_3}^{eff} = -\frac{|J_{i_1}|}{4} + \frac{1}{2} J_{i_0} \vec{S}_{i_0} \vec{S}'_{i_1} + \frac{1}{2} J_{i_2} \vec{S}'_{i_1} \vec{S}_{i_3} \quad (2.29)$$

Rule 2) **AF bond between two spin-1/2**

Here, we directly apply the Ma-Dasgupta procedure discussed in II A: we remove the two spin-1/2 \vec{S}_{i_1} and \vec{S}_{i_2} and replace H_{i_0, i_1, i_2, i_3} by

$$H_{i_0, i_3}^{eff} = -\frac{3}{4} J_{i_1} - \frac{3}{16 J_{i_1}} (J_{i_0}^2 + J_{i_2}^2) + \frac{J_{i_0} J_{i_2}}{2 J_{i_1}} \vec{S}_{i_0} \cdot \vec{S}_{i_3} \quad (2.30)$$

Rule 3) **AF bond between one spin-1 and one spin-1/2**

Suppose that $s_{i_1} = 1$ and $s_{i_2} = \frac{1}{2}$. The hamiltonian $h_0 = J_{i_1} \vec{S}_{i_1} \vec{S}_{i_2}$ admits two energy-levels : the doublet $e_{1/2} = -J_{i_1}$ and the quadruplet $e_{3/2} = \frac{J_{i_1}}{2}$. At first order perturbation theory, Wigner-Eckart theorem gives that, within the subspace of the doublet $s = \frac{1}{2}$, the perturbation h_1 is equivalent to

$$h_1^{eq} = \left(\alpha_1 J_{i_0} \vec{S}_{i_0} + \alpha_2 J_{i_2} \vec{S}_{i_3} \right) \cdot (\vec{S}_{i_1} + \vec{S}_{i_2}) \quad (2.31)$$

where the constants α_1 and α_2 read

$$\alpha_1 = \frac{1}{2} \left[1 + \frac{s_{i_1}(s_{i_1} + 1) - s_{i_2}(s_{i_2} + 1)}{s(s + 1)} \right] = \frac{4}{3} \quad \text{and} \quad \alpha_2 = 1 - \alpha_1 = -\frac{1}{3} \quad (2.32)$$

The renormalization rule is therefore the following : we eliminate the spins \vec{S}_{i_1} and \vec{S}_{i_2} , and replace them by a single spin-1/2 \vec{S}'_{i_1} , and we replace H_{i_0, i_1, i_2, i_3} by the effective Hamiltonian

$$H_{i_0, i'_1, i_3}^{eff} = -J_{i_1} + \frac{4}{3} J_{i_0} \vec{S}_{i_0} \vec{S}'_{i_1} - \frac{1}{3} J_{i_2} \vec{S}'_{i_1} \vec{S}_{i_3} \quad (2.33)$$

Rule 4) AF bond between two spin-1

In this case we apply the rule explained at the beginning of this section (see eqs (2.14)-(2.18)) : we replace the two spin-1 \vec{S}_1 and \vec{S}_2 by two spin-1/2 \vec{S}'_{i_1} and \vec{S}'_{i_2} , and we replace $H_{0,1,2,3}$ by an effective Hamiltonian

$$H_{i_0,i_1,i_2,i_3}^{eff} = -\frac{5J_{i_1}}{4} + J_{i_0}\vec{S}_{i_0}.\vec{S}'_{i_1} + J_{i_1}\vec{S}'_{i_1}.\vec{S}'_{i_2} + J_{i_2}\vec{S}'_{i_2}.\vec{S}_{i_3}. \quad (2.34)$$

This renormalization procedure is entirely consistent from the point of view of the progressive elimination of the highest energy degrees of freedom : it is easy to show that in the four cases of renormalization of a bond described above, all the energy scales Δ_i of the new bonds are always smaller than the energy scale Δ_{i1} of the bond that we renormalize.

It is also easy to check that this renormalization procedure is “closed” inside the set of spin chains defined by the condition (2.21) : if we apply this procedure to an initial chain belonging to this space, such as the random antiferromagnetic spin-1 chain we are interested in, the effective chain always belongs to this set of spin chains. In particular, spins higher than 1 cannot appear through this renormalization scheme.

However, since this renormalization procedure is not purely based on complete decimation of bonds, it introduces correlations between bonds, so that it is impossible to write exact closed flow equations for the probability distributions of couplings, in contrast with the Ma-Dasgupta procedure. To study the properties of this renormalization scheme, we have therefore performed numerical simulations on spin-1 chains containing N sites with periodic boundary conditions ($N = 2^{22}$ for example), whose initial couplings J_i are distributed according to probability distributions of the following form

$$P_d(J) = \frac{1}{d} \text{ for } 1 \leq J \leq 1 + d \quad , \quad \text{and } P_d(J) = 0 \text{ elsewhere} \quad (2.35)$$

The parameter d represents the strength of the initial disorder of the couplings. For a given number of sites N , and a given initial strength d of the disorder, we have numerically implemented the renormalization rules on a given number (typically 100) of initial independent samples, to compute averaged quantities over these different realizations of the initial disorder. It is convenient to use the variable :

$$\Gamma = \ln \frac{\Omega_0}{\Omega}, \quad (2.36)$$

where Ω is the current strongest Δ (see eq. 2.26) and Ω_0 the initial strongest Δ . We have studied the flow of the following quantities : the number $N(\Gamma)$ of effective spins $S = 1/2$ and $S = 1$ still present at scale Γ ; the proportion $\{N_{(S=1)(\Gamma)}/N(\Gamma)\}$ of spins $S = 1$ among the effective spins at scale Γ ; the proportions $\rho_i(\Gamma) = \{N_i(\Gamma)/N(\Gamma)\}$ of bonds of type $i = 1, 2, 3, 4$ at scale Γ ; the probability distributions $P_i(J, \Omega)$ of the coupling J at scale Ω for the four types of bonds $i = 1, 2, 3, 4$. It is in fact more convenient to study the probability distributions $\mathcal{P}_i(x, \Gamma)$ of the reduced variable

$$x = \ln \left(\frac{\Omega}{\Delta(J)} \right), \quad (2.37)$$

where $\Delta(J)$ is defined as in (2.26)

$$\Delta(J) = -J \quad \text{for bonds of type 1} \quad (2.38)$$

$$\Delta(J) = J \quad \text{for bonds of type 2} \quad (2.39)$$

$$\Delta(J) = \frac{3}{2}J \quad \text{for bonds of type 3} \quad (2.40)$$

$$\Delta(J) = 3J \quad \text{for bonds of type 4} \quad (2.41)$$

so that the random variable x varies in $(0, \infty)$ for any type of bonds.

III. THE WEAK DISORDER PHASE

A. Numerical results

In the weak disorder phase, we find that the number $N(\Gamma)$ of effective spins decays exponentially (see Fig 1)

$$N(\Gamma) \underset{\Gamma \rightarrow \infty}{\propto} e^{-\alpha(d)\Gamma} \quad (3.1)$$

where $\alpha(d)$ is a decreasing function of the disorder d that vanishes in the limit $d \rightarrow d_c^-$. As a consequence the magnetic susceptibility at temperature T can be computed by summing Curie laws for the free spins at scale $\Omega = T$. So we have :

$$\chi \propto \frac{1}{T^{1-\alpha(d)}}. \quad (3.2)$$

The proportions $\rho_i(\Gamma)$ of the four types of bonds reach a stationary regime characterized by (see Fig 2)

$$\rho_1(\Gamma) \simeq 0.25 \quad \rho_2(\Gamma) \simeq 0.75 \quad \rho_3(\Gamma) \simeq 0 \quad \rho_4(\Gamma) \simeq 0 \quad (3.3)$$

There are asymptotically only bonds of type 1 and bonds of type 2. This means in particular that there are only effective spin-1/2 in the chain, and no more spin-1. Since two bonds of type 1 cannot be neighbors according to the constraint (2.21), the even bonds and the odds bonds are not equivalent, as in the effective model of Hyman and Yang [9] : the “even” bonds are all antiferromagnetic, whereas the “odd” bonds are either ferromagnetic or antiferromagnetic with equal probability.

It is necessary to introduce the probability distribution $\mathcal{P}_2^{even}(x, \Gamma)$ for the couplings of the even bonds of type 2, and the probability distribution $\mathcal{P}_2^{odd}(x, \Gamma)$ for the couplings of the odd bonds of type 2. We find that $\mathcal{P}_2^{even}(x, \Gamma)$ becomes stationary for large enough Γ , and takes the form of an exponential distribution

$$\mathcal{P}_2^{even}(x, \Gamma) \simeq \alpha_e e^{-\alpha_e x} \quad (3.4)$$

where α_e is independent of Γ , but depends on the value d of the disorder, and is numerically very close to the parameter $\alpha(d)$ characterizing the decay of $N(\Gamma)$ (3.1). The probability

distributions $\mathcal{P}_1(x, \Gamma)$ and $\mathcal{P}_2^{odd}(x, \Gamma)$ coincide (up to statistical fluctuations) and take the form of an exponential distribution (see Fig. 3)

$$\mathcal{P}_1(x, \Gamma) \simeq \mathcal{P}_2^{odd}(x, \Gamma) \simeq \alpha_o(\Gamma) e^{-\alpha_o(\Gamma)x} \quad (3.5)$$

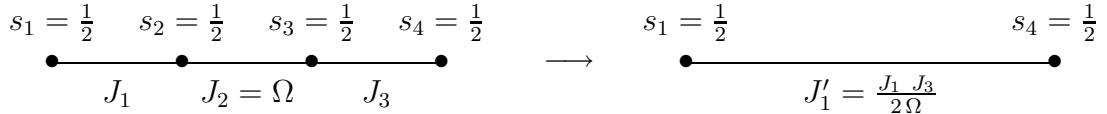
where the parameter $\alpha_o(\Gamma)$ decays exponentially

$$\alpha_o(\Gamma) \propto e^{-\alpha_e \Gamma} \quad (3.6)$$

As a consequence, for large enough Γ , the bond of the chain of highest Δ (corresponding to smallest x) that is chosen to be renormalized, is always an even bond of type 2. In the renormalization operation (2), this even bond disappear together with its two odd neighbors, and a new weak odd bond is produced. This explains why the distribution $\mathcal{P}_2^{even}(x, \Gamma)$ for even bonds remains stationary, whereas the distribution of couplings of odd bonds becomes broader and broader in the variable x . This weak disorder phase is therefore the same as the “Haldane phase” found by Hyman and Yang in their effective model introduced in [9], and is very similar to the random dimer phase found in the study of random dimerized antiferromagnetic spin-1/2 chains [11] : in the asymptotic regime, the chain is made of a set of nearly uncoupled dimers.

B. Approximate flow equations

Assuming that the “even” bonds are all of type 2, that the “odd” bonds are either of type 1 or of type 2 with equal probability, and that the unique important process is the decimation of an even bond according to the rule (2)



it is possible to write approximate flow equations for the probability distributions of the couplings are normalized according to

$$1 = \int_0^\Omega dJ P_2^{even}(J, \Omega) = \int_0^\Omega dJ P_2^{odd}(J, \Omega) = \int_{-\Omega}^0 dJ P_1(J, \Omega) \quad (3.7)$$

It is convenient to introduce the normalized distribution of all odd bonds

$$P^{odd}(J, \Omega) = \frac{1}{2} (P_2^{odd}(J, \Omega) + P_1(J, \Omega)) \quad \text{for } -\Omega < J < \Omega \quad (3.8)$$

The approximate flow equations for the probability distributions $P_2^{even}(J, \Omega)$ and $P_2^{odd}(J, \Omega)$ then read

$$-\frac{\partial P_2^{even}(J, \Omega)}{\partial \Omega} = P_2^{even}(\Omega, \Omega) P_2^{even}(J, \Omega) \quad (3.9)$$

$$-\frac{\partial P^{odd}(J, \Omega)}{\partial \Omega} = -P_2^{even}(\Omega, \Omega) P^{odd}(J, \Omega) \quad (3.10)$$

$$+ P_2^{even}(\Omega, \Omega) \int_{-\Omega}^{\Omega} dJ_1 P^{odd}(J_1, \Omega) \int_{-\Omega}^{\Omega} dJ_3 P^{odd}(J_3, \Omega) \delta\left(J - \frac{J_1 J_3}{2\Omega}\right) \quad (3.11)$$

In the new variables $\Gamma = \ln \frac{\Omega_0}{\Omega}$ and $x = \ln \left(\frac{\Omega}{J}\right) \in [0, +\infty)$, the flow equation for $\mathcal{P}_2^{even}(x)$ admits stationary solutions of exponential form

$$\mathcal{P}_2^{even}(x) = \alpha_e e^{-\alpha_e x} \quad (3.12)$$

with undetermined constant α_e , in agreement with our numerical result (3.4). With the last change of variables

$$x \longrightarrow z = \alpha_o(\Gamma) \ln \left(\frac{\Omega}{|J|} \right) \quad (3.13)$$

the flow equation for the corresponding probability distributions $\tilde{P}_1^{odd}(z, \Gamma)$ and $\tilde{P}_2^{odd}(z, \Gamma)$ admit the same stationary solution

$$\tilde{P}_1^{odd}(z, \Gamma) \simeq \tilde{P}_2^{odd}(z, \Gamma) \xrightarrow{\Gamma \rightarrow \infty} e^{-z} \quad \text{with } \alpha_o(\Gamma) \xrightarrow{\Gamma \rightarrow \infty} e^{-\alpha_e \Gamma} \quad (3.14)$$

where α_e is the number characterizing $\mathcal{P}_2^{even}(x)$ (3.12). We may also write the flow equation for the total number $N(\Omega)$ of spins still present at scale Ω

$$-\frac{dN}{d\Omega} = -P_2^{even}(\Omega, \Omega) N(\Omega), \quad (3.15)$$

so that we obtain the following asymptotic behavior in the variable Γ :

$$N(\Gamma) \xrightarrow{\Gamma \rightarrow \infty} e^{-\alpha_e \Gamma}. \quad (3.16)$$

IV. THE STRONG DISORDER PHASE

A. Numerical results

In the strong disorder phase $d > d_c$, we find that the number $N(\Gamma)$ of effective spins decays as in the random singlet theory for the disordered antiferromagnetic spin-1/2 chain (see Fig 4) :

$$N(\Gamma) \xrightarrow{\Gamma \rightarrow \infty} \frac{1}{\Gamma^2}. \quad (4.1)$$

The magnetic susceptibility has thus the random singlet behaviour :

$$\chi \propto \frac{1}{T \log^2 T}. \quad (4.2)$$

The proportions $\rho_i(\Gamma)$ of the four types of bonds reach an asymptotic regime characterized by (see Fig 5)

$$\rho_1(\Gamma) \sim 0 \quad \rho_2(\Gamma) \sim \epsilon(\Gamma) \quad \rho_3(\Gamma) \sim 2\epsilon(\Gamma) \quad \rho_4(\Gamma) \sim 1 - 3\epsilon(\Gamma) \quad (4.3)$$

where $\epsilon(\Gamma)$ slowly goes to 0 as $\Gamma \rightarrow \infty$. This means that there is a sea of bonds of type 4, with sometimes defects of structure { bond of type 3, bond of type 2, bond of type 3 }. This defect structure is produced by the renormalization rule 4) for a bond of type 4 when its two neighbor bonds are also of type 4. The fact that there is no more bonds of type 1 in the asymptotic regime (4.3) shows that defects are destroyed by the renormalization of the central bond of type 2 and not by the bonds of type 3 ; this means that for the probability distribution $P_4(J, \Omega)$ at large enough Ω , two typical couplings are much weaker than the bigger one. We indeed find that $\mathcal{P}_4(x, \Gamma)$ is an exponential distribution (see Fig 6)

$$\mathcal{P}_4(x, \Gamma) \simeq \alpha_4(\Gamma) e^{-\alpha_4(\Gamma)x} \quad (4.4)$$

where the parameter $1/\alpha_4(\Gamma)$ follows the random singlet behavior (see Fig. 7)

$$\frac{1}{\alpha_4(\Gamma)} \simeq \Gamma + \text{Cst} \quad (4.5)$$

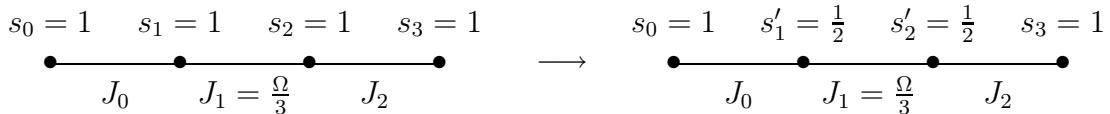
As a consequence, if a defect is produced at the renormalization energy scale Ω , it survives until the energy scale $\frac{\Omega}{3}$ where it get decimated according to the rule (2), and the whole defect of structure { bond of type 3, bond of type 2, bond of type 3 } entirely disappears to give one bond of type 4. Fig 8 shows indeed clearly that the probability distribution $\mathcal{P}_2(x, \Gamma)$ tends to concentrate on the interval $0 < x < \ln 3$ as Γ increases. That has to be contrasted with the bonds of type 3, which are characterized by a distribution $\mathcal{P}_3(x, \Gamma)$ that tends to coincide with $\mathcal{P}_4(x, \Gamma)$ for large enough Γ .

B. Approximate flow equations phase

Assuming that there is a sea of bonds of type 4, with sometimes defects of structure { bond of type 3, bond of type 2, bond of type 3 }, it is possible to write approximate flow equations for the probability distributions of the couplings normalized according to :

$$1 = \int_0^{\Omega} dJ P_2(J, \Omega) = \int_0^{\frac{2\Omega}{3}} dJ P_3(J, \Omega) = \int_0^{\frac{\Omega}{3}} dJ P_4(J, \Omega). \quad (4.6)$$

Assuming that the only two important renormalization processes are the production of the defect structure bond of type 3, bond of type 2, bond of type 3 by the renormalization rule 4) for a bond of type 4 when its two neighbor bonds are also of type 4



and the suppression of the defect structure by the decimation rule 2)



we obtain the following approximate flow equations for the three probability distributions :

$$-\frac{\partial P_2(J, \Omega)}{\partial \Omega} = P_2(\Omega, \Omega) P_2(J, \Omega) + \frac{1}{3} P_4\left(\frac{\Omega}{3}, \Omega\right) \frac{N_4(\Omega)}{N_2(\Omega)} \left[\delta\left(J - \frac{\Omega}{3}\right) - P_2(J, \Omega) \right] \quad (4.7)$$

$$-\frac{\partial P_3(J, \Omega)}{\partial \Omega} = \frac{2}{3} P_4\left(\frac{\Omega}{3}, \Omega\right) \frac{N_4(\Omega)}{N_3(\Omega)} \left[P_4(J, \Omega) - P_3(J, \Omega) \right] \quad (4.8)$$

$$\begin{aligned} -\frac{\partial P_4(J, \Omega)}{\partial \Omega} &= \frac{1}{3} P_4\left(\frac{\Omega}{3}, \Omega\right) P_4(J, \Omega) - \frac{N_2(\Omega)}{N_4(\Omega)} P_2(\Omega, \Omega) P_4(J, \Omega) \\ &+ \frac{N_2(\Omega)}{N_4(\Omega)} P_2(\Omega, \Omega) \int_0^{\frac{2\Omega}{3}} dJ_0 P_3(J_0, \Omega) \int_0^{\frac{2\Omega}{3}} dJ_2 P_3(J_2, \Omega) \delta\left(J - \frac{J_0 J_2}{2\Omega}\right), \end{aligned} \quad (4.9) \quad (4.10)$$

together with the flow equations for the number $N_i(\Omega)$ of bonds of type $i = 2, 3, 4$:

$$-\frac{dN_2}{d\Omega} = -\frac{1}{2} \frac{dN_3}{d\Omega} = \frac{1}{3} P_4\left(\frac{\Omega}{3}, \Omega\right) N_4(\Omega) - P_2(\Omega, \Omega) N_2(\Omega) \quad (4.11)$$

$$-\frac{dN_4}{d\Omega} = P_2(\Omega, \Omega) N_2(\Omega) - P_4\left(\frac{\Omega}{3}, \Omega\right) N_4(\Omega), \quad (4.12)$$

so that the total number $N(\Omega) = N_2(\Omega) + N_3(\Omega) + N_4(\Omega)$ of bonds evolves according to :

$$-\frac{dN}{d\Omega} = -2P_2(\Omega, \Omega) N_2(\Omega). \quad (4.13)$$

It is more convenient to write the flow equations for the probability distributions $\mathcal{P}_i(x, \Gamma)$ of the reduced variable $x = \ln\left(\frac{\Omega}{\Delta(J)}\right)$, where $\Delta(J)$ is defined by (2.41) so that the random variable x varies in $(0, \infty)$ for any type of bonds :

$$\begin{aligned} \frac{\partial \mathcal{P}_2(x, \Gamma)}{\partial \Gamma} &= \frac{\partial \mathcal{P}_2(x, \Gamma)}{\partial x} + \mathcal{P}_2(0, \Gamma) \mathcal{P}_2(x, \Gamma) + \frac{N_4(\Gamma)}{N_2(\Gamma)} \mathcal{P}_4(0, \Gamma) \left[\delta(x - \ln 3) - \mathcal{P}_2(x, \Gamma) \right], \\ & \end{aligned} \quad (4.14)$$

$$\frac{\partial \mathcal{P}_3(x, \Gamma)}{\partial \Gamma} = \frac{\partial \mathcal{P}_3(x, \Gamma)}{\partial x} + 2 \frac{N_4(\Gamma)}{N_2(\Gamma)} \mathcal{P}_4(0, \Gamma) \left[\mathcal{P}_4(x - \ln 2, \Gamma) - \mathcal{P}_3(x, \Gamma) \right], \quad (4.15)$$

$$\frac{\partial \mathcal{P}_4(x, \Gamma)}{\partial \Gamma} = \frac{\partial \mathcal{P}_4(x, \Gamma)}{\partial x} + \left[\mathcal{P}_4(0, \Gamma) - \frac{N_2(\Gamma)}{N_4(\Gamma)} \mathcal{P}_2(0, \Gamma) \right] \mathcal{P}_4(x, \Gamma) \quad (4.16)$$

$$+ \frac{N_2(\Gamma)}{N_4(\Gamma)} \mathcal{P}_2(0, \Gamma) \int_0^\infty dx_1 \mathcal{P}_3(x_1, \Gamma) \int_0^\infty dx_2 \mathcal{P}_3(x_2, \Gamma) \delta \left(x - x_1 - x_2 - \ln \frac{3}{2} \right). \quad (4.17)$$

Since the singular term containing the delta-function in (4.14) tends to develop a discontinuity in $\mathcal{P}_2(x, \Gamma)$ at $x = \ln 3$, it is convenient to set :

$$\mathcal{P}_2(x, \Gamma) = [1 - e(\Gamma)] \frac{\theta(\ln 3 - x)}{\ln 3} + e(\Gamma) f_2(x, \Gamma), \quad (4.18)$$

where $0 < e(\Gamma) < 1$ and $f_2(x, \Gamma)$ is a normalized probability distribution that is regular at $x = \ln 3$. Equation (4.14) will be satisfied if $e(\Gamma)$ and $f(x, \Gamma)$ satisfy :

$$e(\Gamma) = 1 - (\ln 3) \frac{N_4(\Gamma)}{N_2(\Gamma)} \mathcal{P}_4(0, \Gamma) \quad (4.19)$$

$$\frac{de(\Gamma)}{d\Gamma} = -e(\Gamma)[1 - e(\Gamma)] f(0, \Gamma) \quad (4.20)$$

$$\frac{\partial f(x, \Gamma)}{\partial \Gamma} = \frac{\partial f(x, \Gamma)}{\partial x} + f(0, \Gamma) f(x, \Gamma). \quad (4.21)$$

Obvious stationary solutions for $f(x, \Gamma)$ are simple exponentials :

$$f(x, \Gamma) \underset{\Gamma \rightarrow \infty}{\simeq} \alpha_f e^{-\alpha_f x}, \quad (4.22)$$

in which case $e(\Gamma)$ vanishes exponentially :

$$e(\Gamma) \underset{\Gamma \rightarrow \infty}{\propto} e^{-\alpha_f \Gamma}, \quad (4.23)$$

so that $\mathcal{P}_2(x, \Gamma)$ converges towards the stationary solution :

$$\mathcal{P}_2^*(x) = \frac{1}{\ln 3} \theta(\ln 3 - x). \quad (4.24)$$

This corresponds in the original variables to

$$P_2(J, \Omega) = \frac{1}{(\ln 3) J} \quad \text{for } \frac{\Omega}{3} < J < \Omega. \quad (4.25)$$

We also obtain the following equation in the asymptotic regime :

$$N_4(\Gamma) \mathcal{P}_4(0, \Gamma) \simeq N_2(\Gamma) \frac{1}{\ln 3}, \quad (4.26)$$

that we will use now to study the flow equations for $\mathcal{P}_3(x, \Gamma)$ and $\mathcal{P}_4(x, \Gamma)$

With the last change of variables :

$$x \longrightarrow z = \alpha_4(\Gamma)x, \quad (4.27)$$

we find that the flow equation for the corresponding probability distributions $\tilde{P}_4(z, \Gamma)$ and $\tilde{P}_3(z, \Gamma)$ admit the stationary solutions :

$$\tilde{P}_4(z, \Gamma) \xrightarrow[\Gamma \rightarrow \infty]{} e^{-z} \quad \text{and} \quad \tilde{P}_3(z, \Gamma) \xrightarrow[\Gamma \rightarrow \infty]{} e^{-z}, \quad (4.28)$$

where

$$\alpha_4(\Gamma) \xrightarrow[\Gamma \rightarrow \infty]{} \frac{1}{\Gamma}, \quad (4.29)$$

as in the random singlet solution of Ma-Dasgupta. It is then easy to obtain the asymptotic behavior of the total number $N(\Gamma)$ of spins (4.13) :

$$N(\Gamma) \xrightarrow[\Gamma \rightarrow \infty]{} \frac{1}{\Gamma^2}, \quad (4.30)$$

and the asymptotic behavior of the proportion $\epsilon(\Gamma)$ (4.3) of defects from (4.26) :

$$\epsilon(\Gamma) = \frac{N_2(\Gamma)}{N(\Gamma)} \xrightarrow[\Gamma \rightarrow \infty]{} \frac{\ln 3}{\Gamma}. \quad (4.31)$$

V. THE CRITICAL REGIME

On Fig. 9, we have plotted the proportion $\frac{N_{(S=1)}(\Gamma)}{N(\Gamma)}$ of spins $S = 1$ among the effective spins for various values of the disorder this proportion flows towards 0 in the weak disorder phase and to 1 in the strong disorder phase. Between these two attractive values, there is an unstable fixed point at $d_c \simeq 5.75(5)$ where the proportion of spins $S = 1$ among the effective spins remains stationary at the intermediate value 0.315(5). The proportions $\rho_i(\Gamma)$ of the four types of bonds reach a stationary state characterized by (see Fig 10)

$$\rho_1(\Gamma) \sim 0.17, \quad \rho_2(\Gamma) \sim 0.35, \quad \rho_3(\Gamma) \sim 0.33, \quad \rho_4(\Gamma) \sim 0.15. \quad (5.1)$$

We find of course that the four probability distributions $\mathcal{P}_i(x, \Gamma)$ for $i = 1, 2, 3, 4$ coincide up to statistical fluctuations (otherwise, the proportions $\rho_i(\Gamma)$ would not remain stationary) and follow the exponential form (see Fig 11) :

$$\mathcal{P}_i(x, \Gamma) \simeq \alpha_c(\Gamma) e^{-\alpha_c(\Gamma)x}, \quad (5.2)$$

where the parameter $1/\alpha_c(\Gamma)$ (see Fig. 12) follows the behavior of the effective model of Hyman and Yang [9] :

$$\frac{1}{\alpha_c(\Gamma)} \simeq \frac{\Gamma}{2} + \text{Cst.} \quad (5.3)$$

The magnetic susceptibility is given by the effective number of free spins :

$$\chi \propto \frac{1}{T \log^3 T}. \quad (5.4)$$

VI. CONCLUSION

We have introduced a real-space renormalization scheme that allows the study of the spin-1 chain. Within this scheme we obtained a complete characterization of the weak-coupling phase, the critical regime and the strong-disorder phase. In all phases we were able to follow the spin populations and to obtain the probability distributions of the different types of bonds that appear under renormalization. It is only in the weak and strong coupling limit that we were able to obtain approximate analytical flow equations.

The renormalization scheme that we used is an extension of the Ma-Dasgupta idea. These schemes have in common the fact that they are consistent for arbitrarily weak initial disorder. They do not create bonds stronger than the original decimated bond. In the spin-1/2 case, it is believed that this means that there is no critical disorder. In fact, this is suggested by bosonization: most bosonic forms of randomness give rise to relevant operators along the massless line of the pure system when the anisotropy is varied. The simplest assumption [2] is thus that the system flows immediately to the random-singlet phase (there is a region of stability of the spin liquid but this happens only for attractive enough interactions between the Jordan-Wigner fermions i.e. for negative enough anisotropy).

However, this is not the case for the spin-1 chain. Here the Haldane gap is perturbatively insensitive to disorder as naively expected. This is known from bosonization studies of the spin-1/2 two-leg ladder [12] as well as of the anisotropic spin-1 chain [13]. So we may be in a situation with a first critical disorder strength corresponding to the vanishing of Haldane gap but which is unreachable by the real-space scheme. With increasing disorder there is then the second critical disorder strength for which the string order vanishes. This second transition is described by our renormalization scheme which is then asymptotically exact. Conversely, the bosonization methods are unable to follow the flow to strong coupling and thus are unable to describe even the weak-disorder phase captured by the real-space scheme. It may be also that there is nothing like a critical value of the disorder for the vanishing of the Haldane gap, if for example there are states of arbitrarily small energies in the gap as in the case of the Lifshitz tails in the localization problem. It remains to be seen if there is a single theoretical approach that is able to deal all known limiting cases.

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FIGURES

Proportion of surviving spins

for weak initial disorder ($d=0.1, d=0.5, d=1$ and $d=2$)

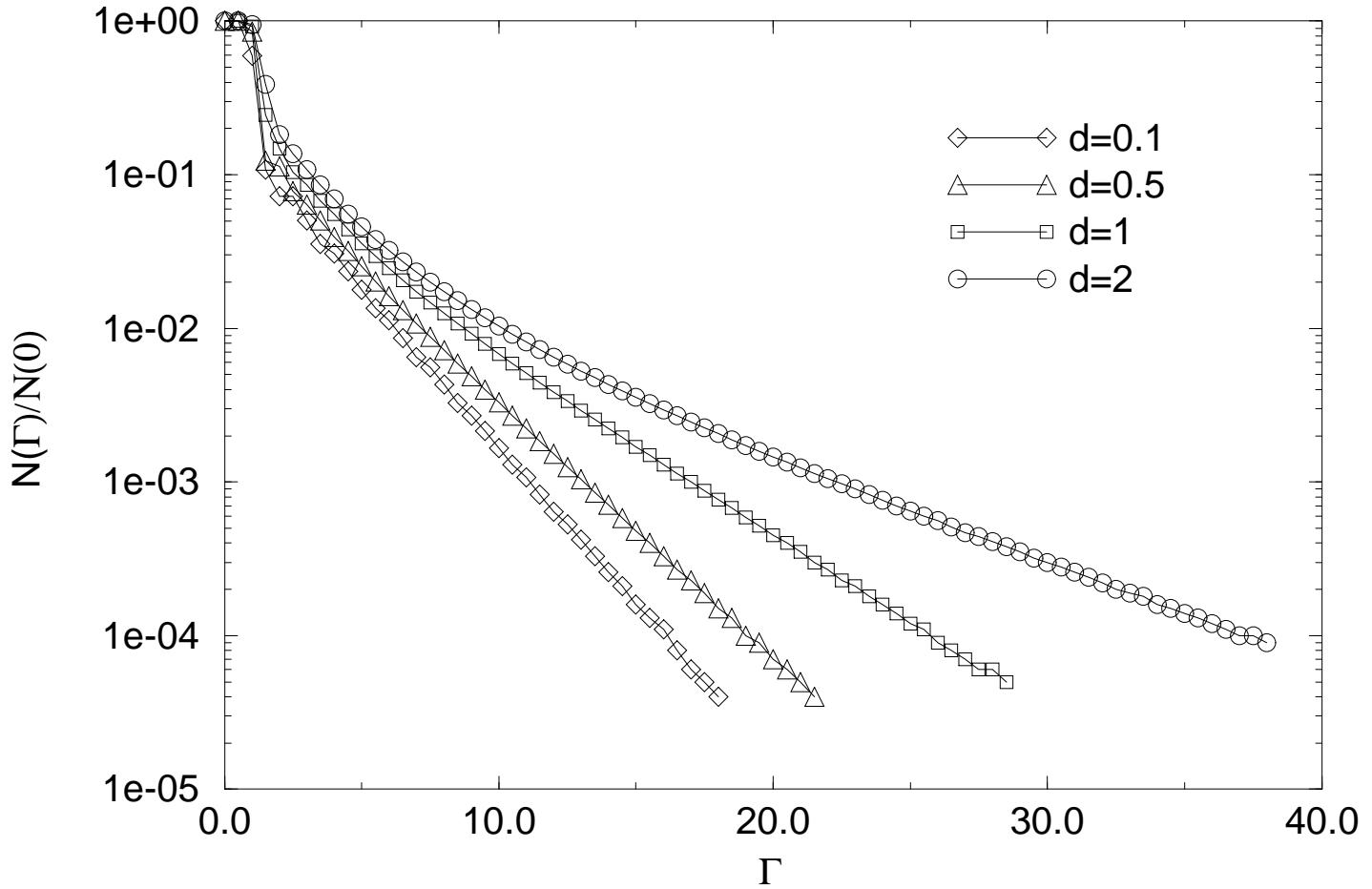


FIG. 1. Linear-Log plot of the proportion $\frac{N(\Gamma)}{N(0)}$ of effective spins at scale Γ , for weak initial disorder $d = 0.1, d = 0.5, d = 1$ and $d = 2$: this proportion decays exponentially (3.1).

Proportions $\rho_1(\Gamma), \rho_2(\Gamma), \rho_3(\Gamma), \rho_4(\Gamma)$

for weak initial disorder ($d=0.1$)

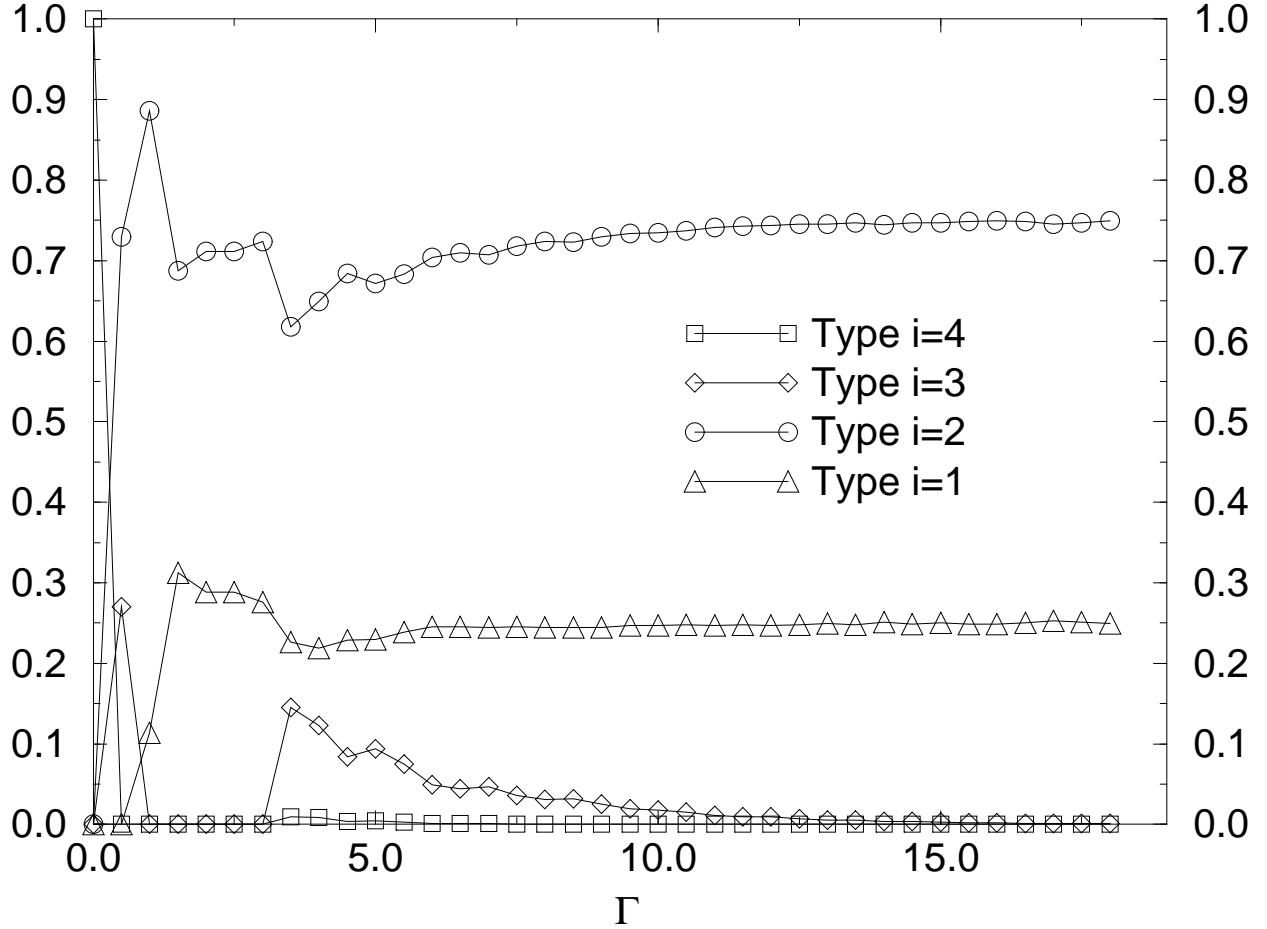


FIG. 2. The proportions $\rho_i(\Gamma)$ of the four types $i = 1, 2, 3, 4$ of bonds at scale Γ , for weak initial disorder $d = 0.1$: they reach the asymptotic regime (3.3).

Probability distribution $P_{2,odd}(x, \Gamma)$

for weak initial disorder $d=0.5$

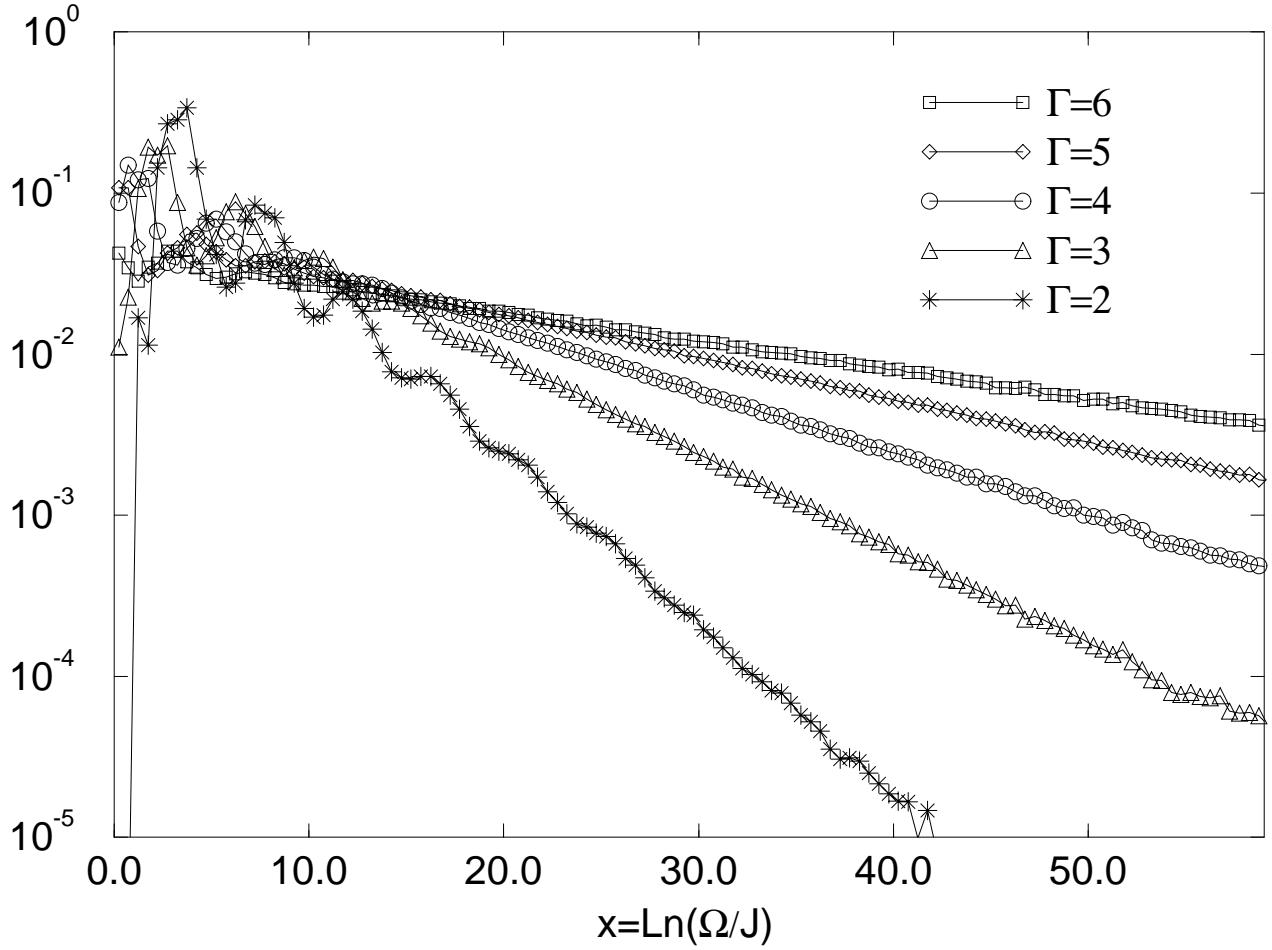


FIG. 3. Linear-Log plot of the probability distribution $P_{2,odd}^{odd}(x, \Gamma)$ for $\Gamma = 2, 3, 4, 5, 6$, for weak initial disorder $d = 0.5$: $P_{2,odd}^{odd}(x, \Gamma)$ is well described by the exponential form (3.5) with a parameter $\alpha_o(\Gamma)$ that is found to decay exponentially with Γ (3.6).

Proportion of surviving spins

for a large initial disorder (d=100)

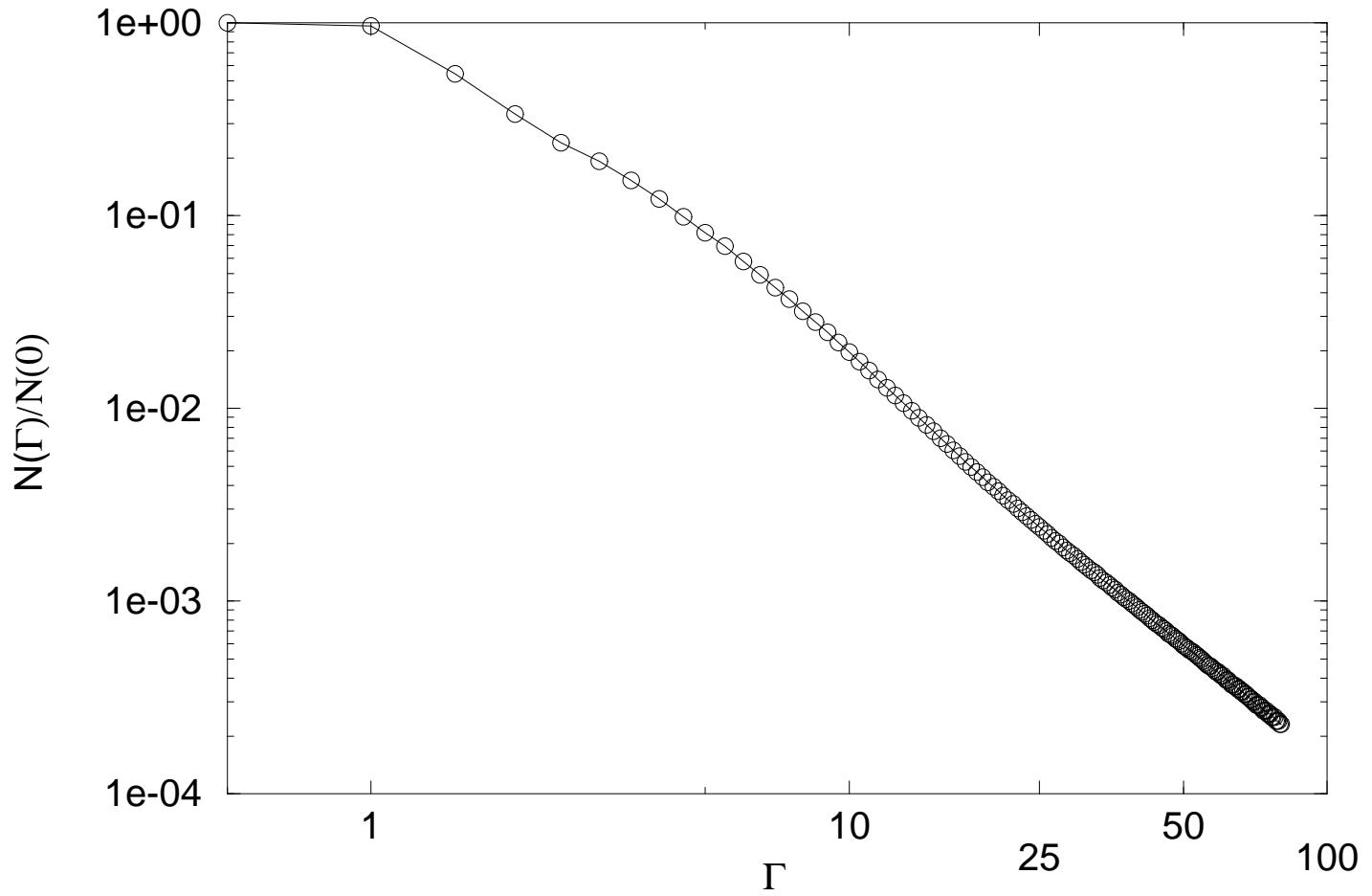


FIG. 4. Log-Log plot of the proportion $\frac{N(\Gamma)}{N(0)}$ of effective spins at scale Γ for strong initial disorder $d = 100$: this proportion follows the power-law asymptotic behavior (4.1).

Proportions $\rho_1(\Gamma), \rho_2(\Gamma), \rho_3(\Gamma)/2, \rho_4(\Gamma)$

for large initial disorder ($d=100$)

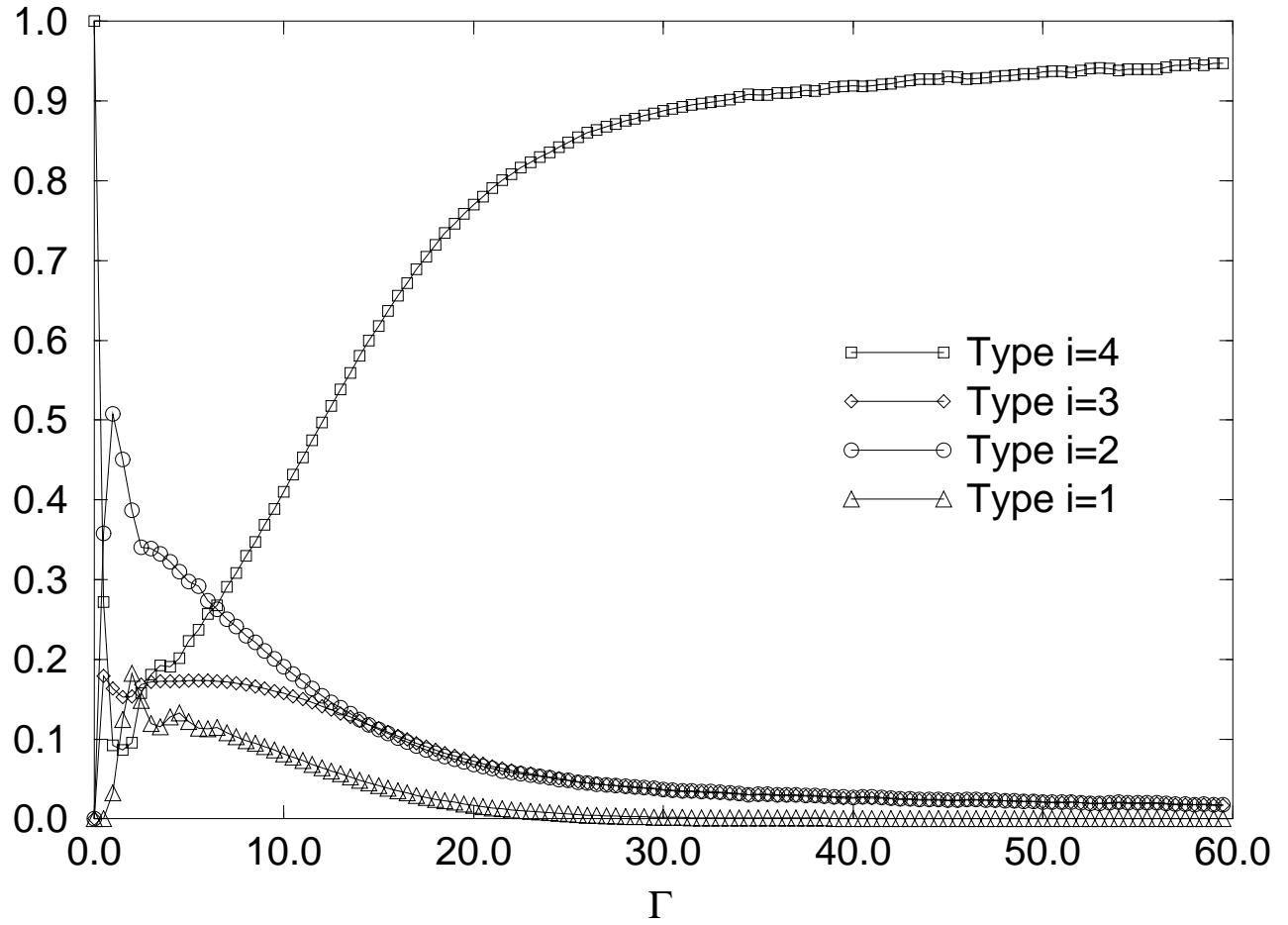


FIG. 5. The proportions $\rho_i(\Gamma)$ of the four types $i = 1, 2, 3, 4$ of bonds at scale Γ , for strong initial disorder $d = 100$: they reach the asymptotic regime (4.3).

Probability distribution $P_4(x, \Gamma)$

for large initial disorder ($d=100$)

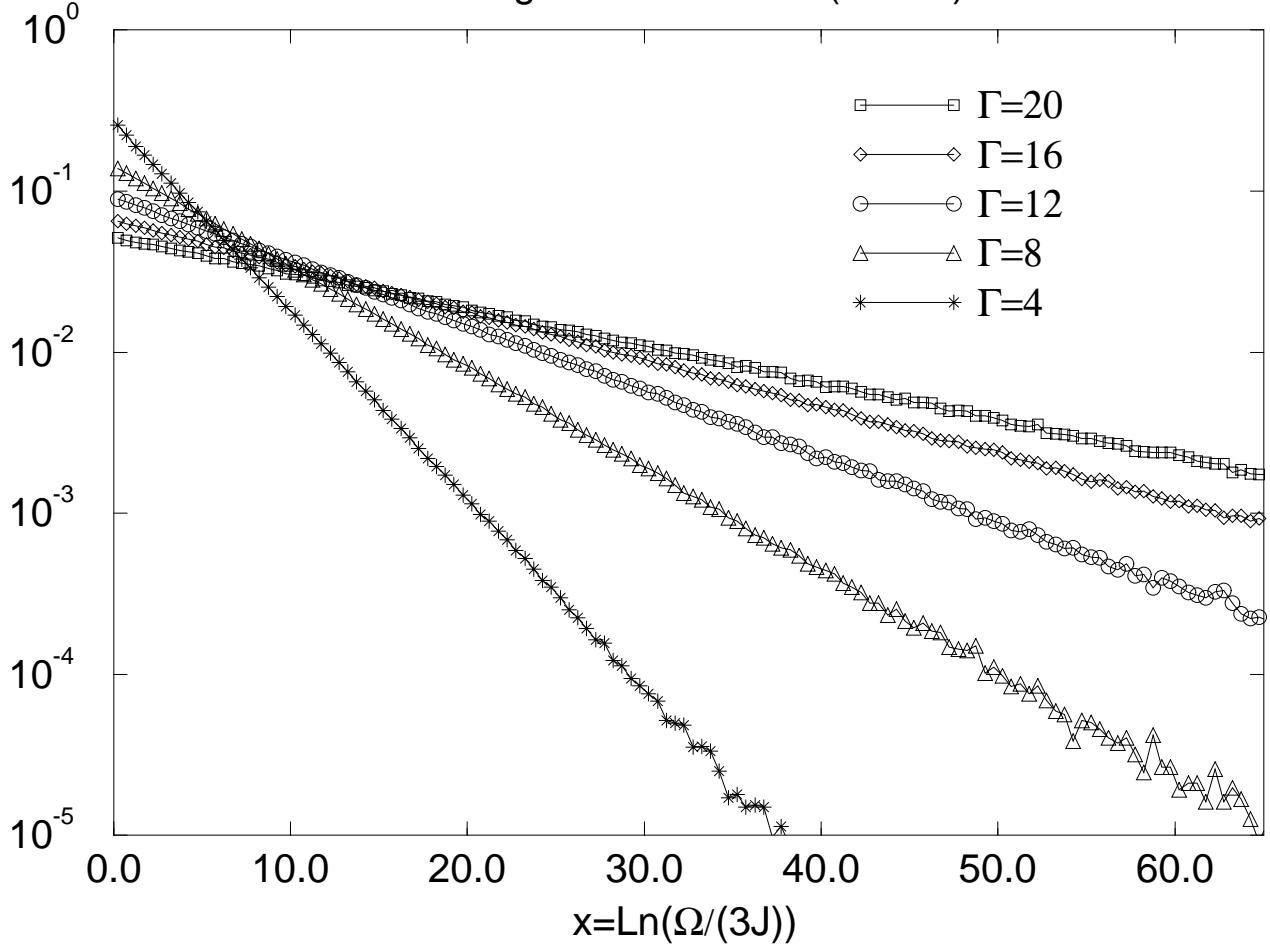


FIG. 6. Linear-Log plot of the probability distribution $P_4(x, \Gamma)$ for $\Gamma = 4, 8, 12, 16, 20$, for strong initial disorder $d = 100$: $P_4(x, \Gamma)$ is well described by the exponential form (4.4) with a parameter $\alpha_4(\Gamma)$ plotted on Fig 7.

$1/\alpha_4(\Gamma)$

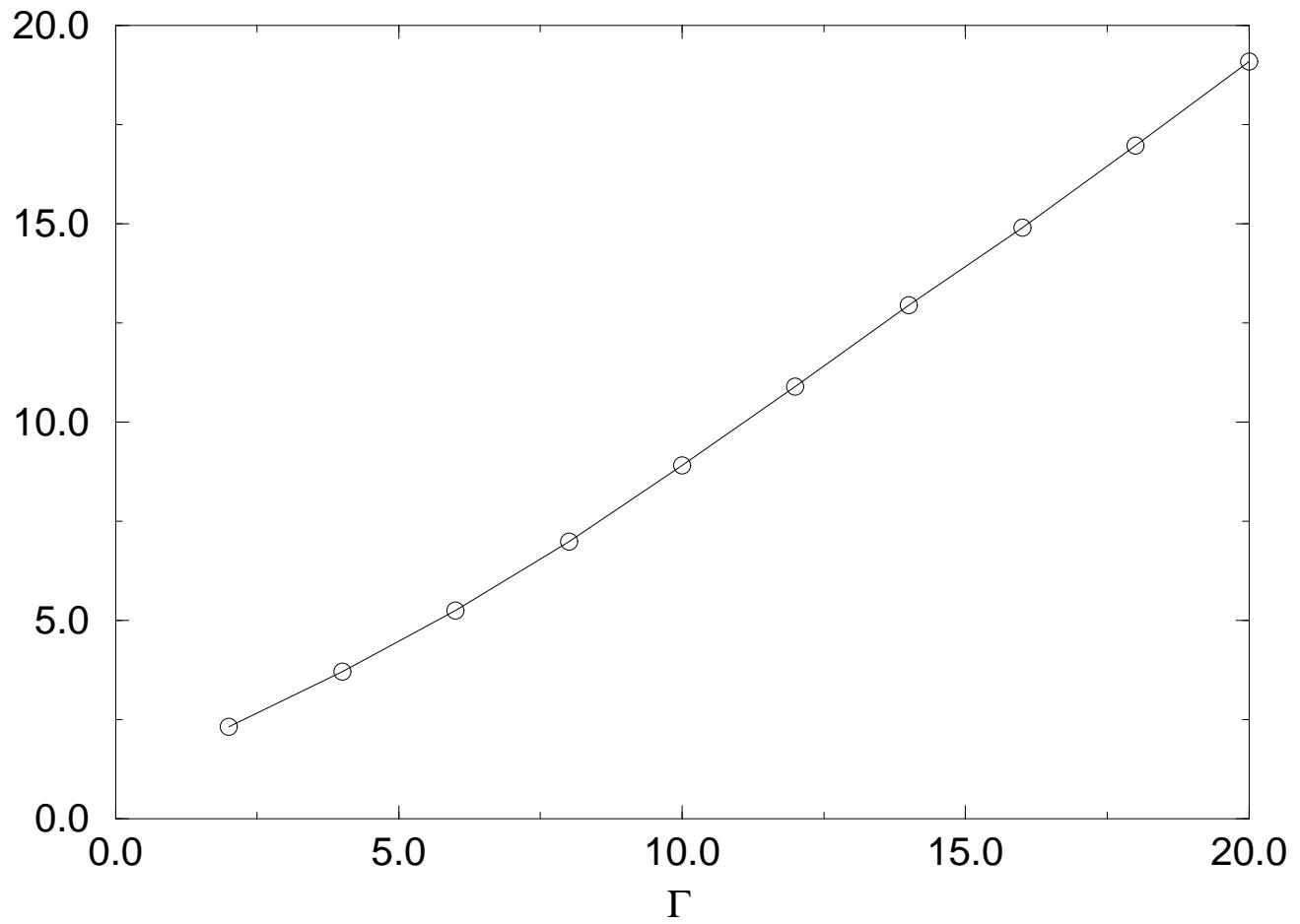


FIG. 7. Plot of the inverse of the parameter $\alpha_4(\Gamma)$ defined in (4.4) as a function of Γ : it follows the random singlet behavior 4.5.

Probability distribution $P_2(x, \Gamma)$

for large initial disorder $d=100$

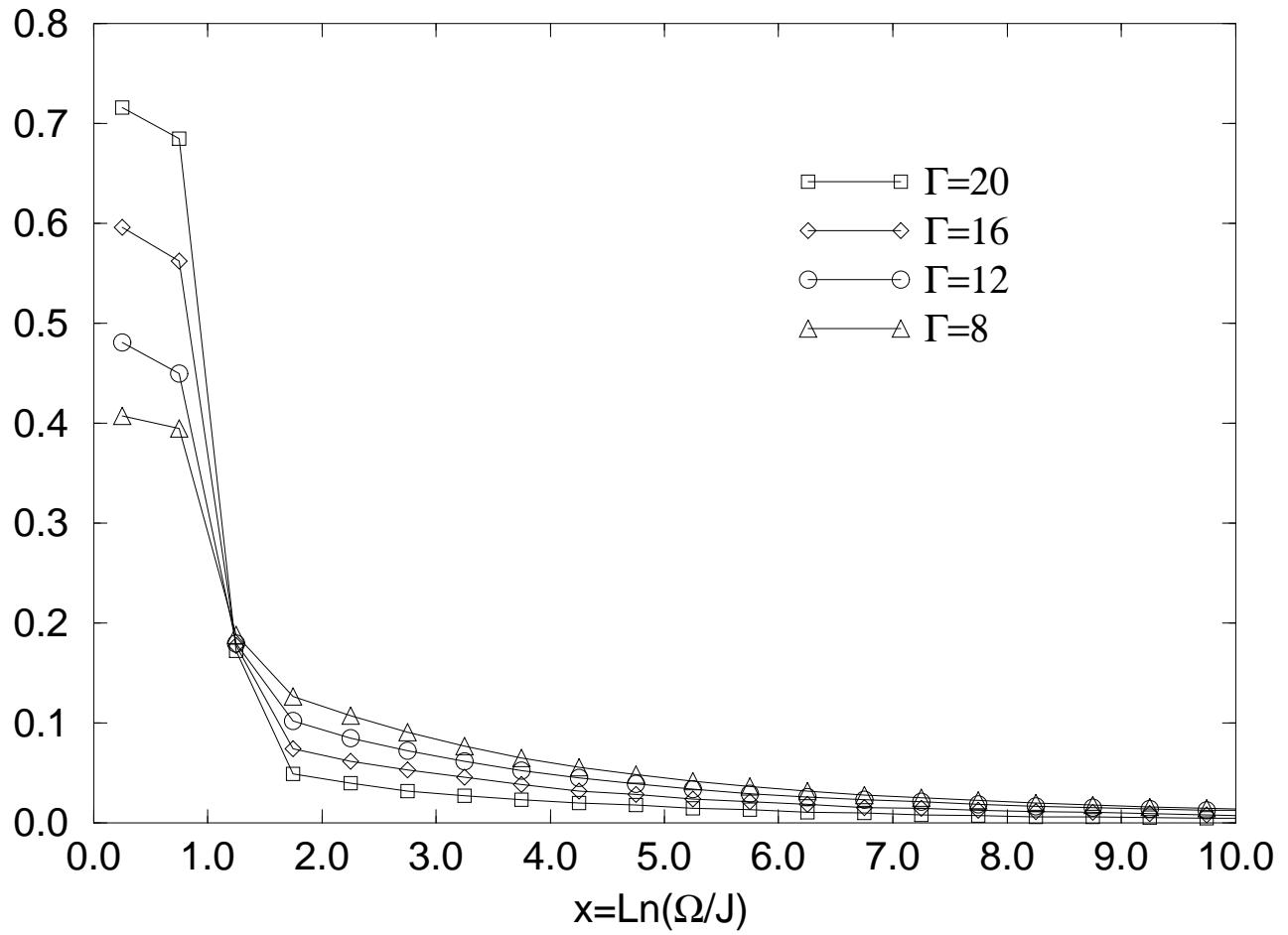


FIG. 8. Plot of the probability distribution $P_2(x, \Gamma)$ for $\Gamma = 8, 12, 16, 20$, for a strong initial disorder $d = 100$: $P_2(x, \Gamma)$ tends to concentrate on the interval $0 < x < \ln 3$ as explained in the text.

Proportion of spin-1 among the effective spins

for initial disorder $d=1, 2, 4, 5.5, 6, 8, 16, 100$

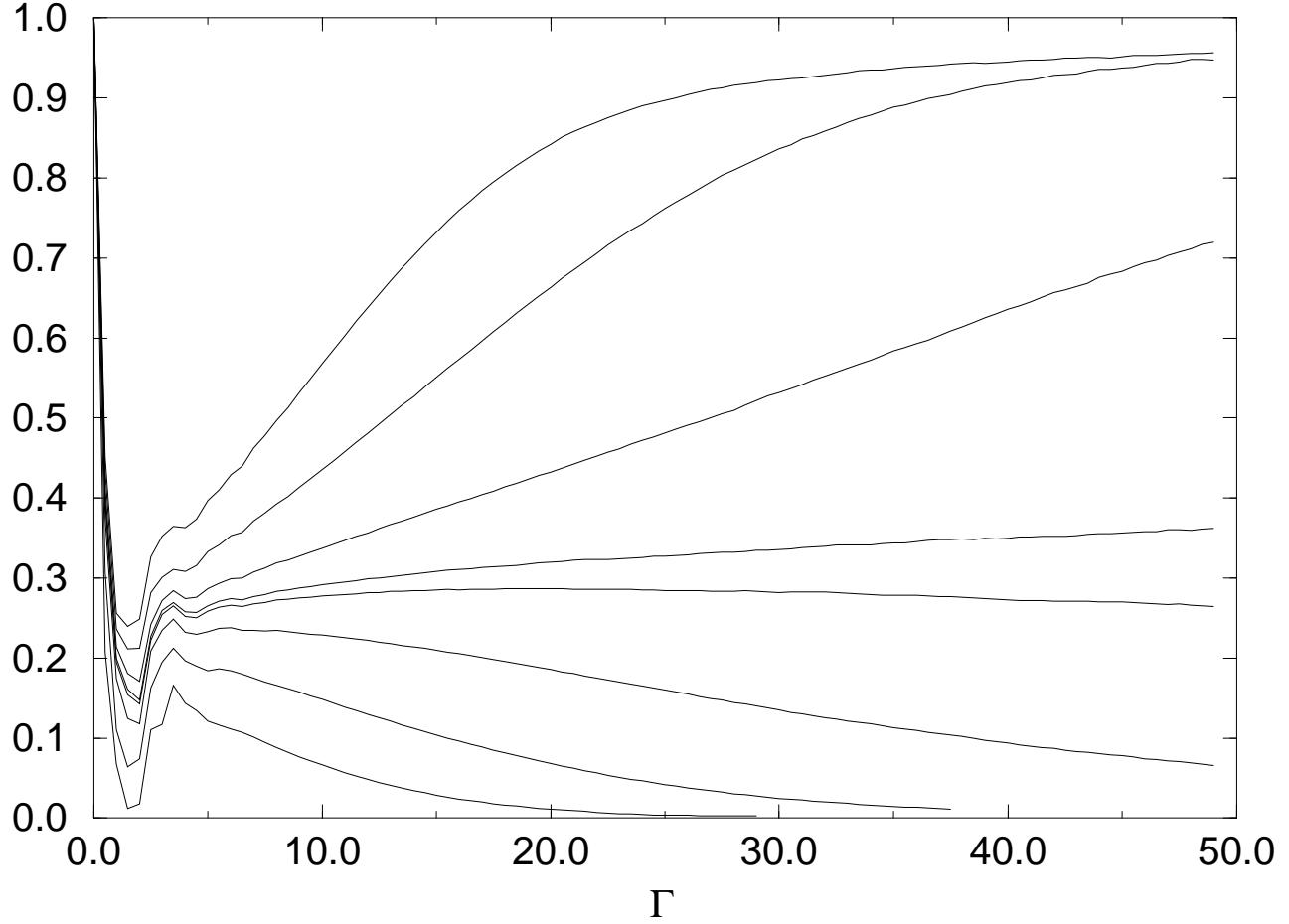


FIG. 9. Proportion of spins $S = 1$ among the effective spins at scale Γ for various values of the initial disorder $d = 1, 2, 3, 4, 5.5, 6, 8, 16, 100$: this proportion flows towards 0 in the weak disorder phase and to 1 in the strong disorder phase. Between these two attractive values, there is an unstable fixed point at $d_c \simeq 5.75(5)$ where the proportion of spins $S = 1$ among the effective spins remains stationary at the intermediate value 0.315(5).

Proportions $\rho_1(\Gamma), \rho_2(\Gamma), \rho_3(\Gamma), \rho_4(\Gamma)$

for critical initial disorder $d_c = 5.75$

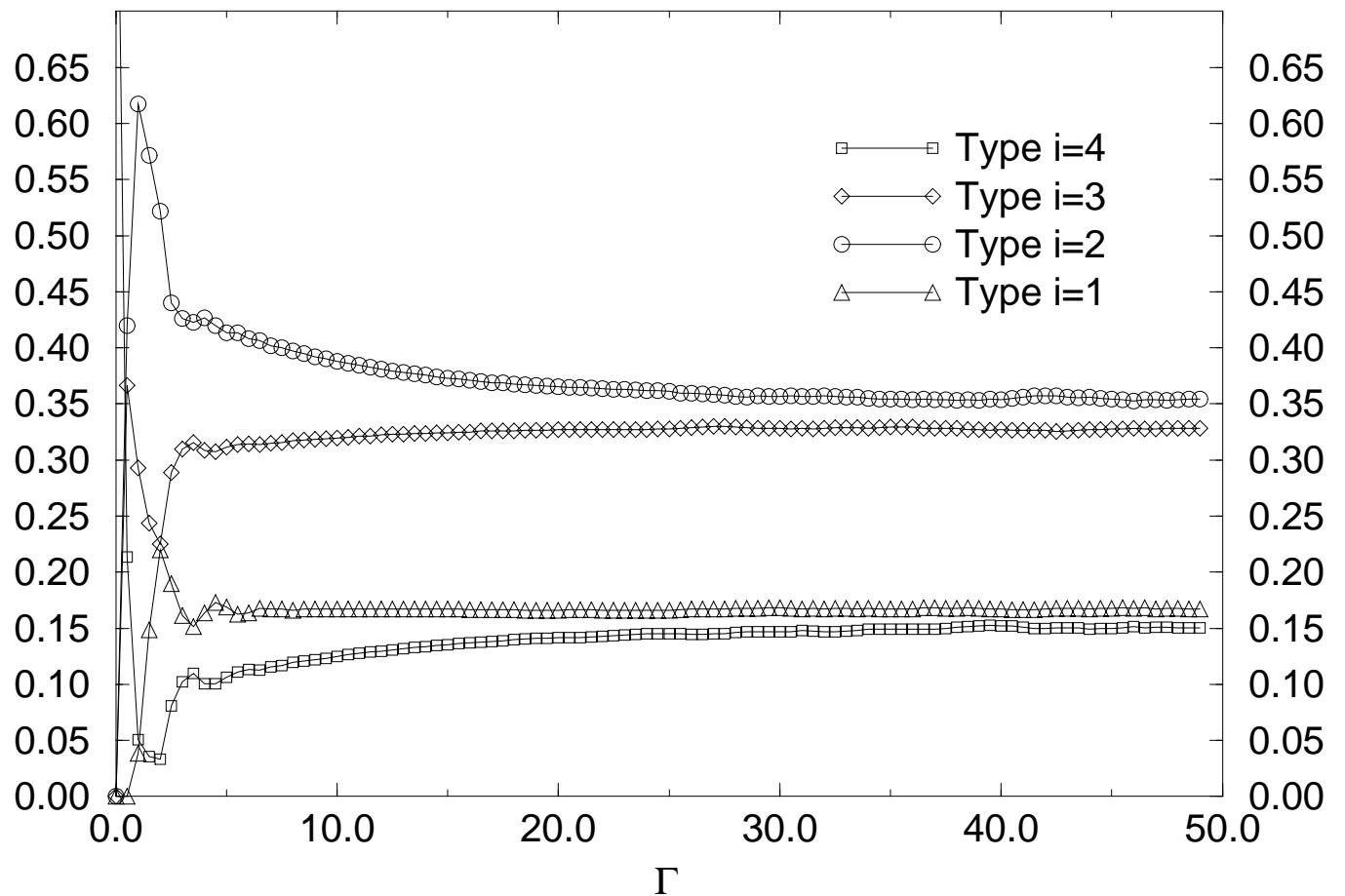


FIG. 10. The proportions $\rho_i(\Gamma)$ of the four types $i = 1, 2, 3, 4$ of bonds at scale Γ , for the critical initial disorder $d_c = 5.75$ (5.1).

Probability distribution $P_2(x, \Gamma)$

for critical initial disorder $d_c = 5.75$

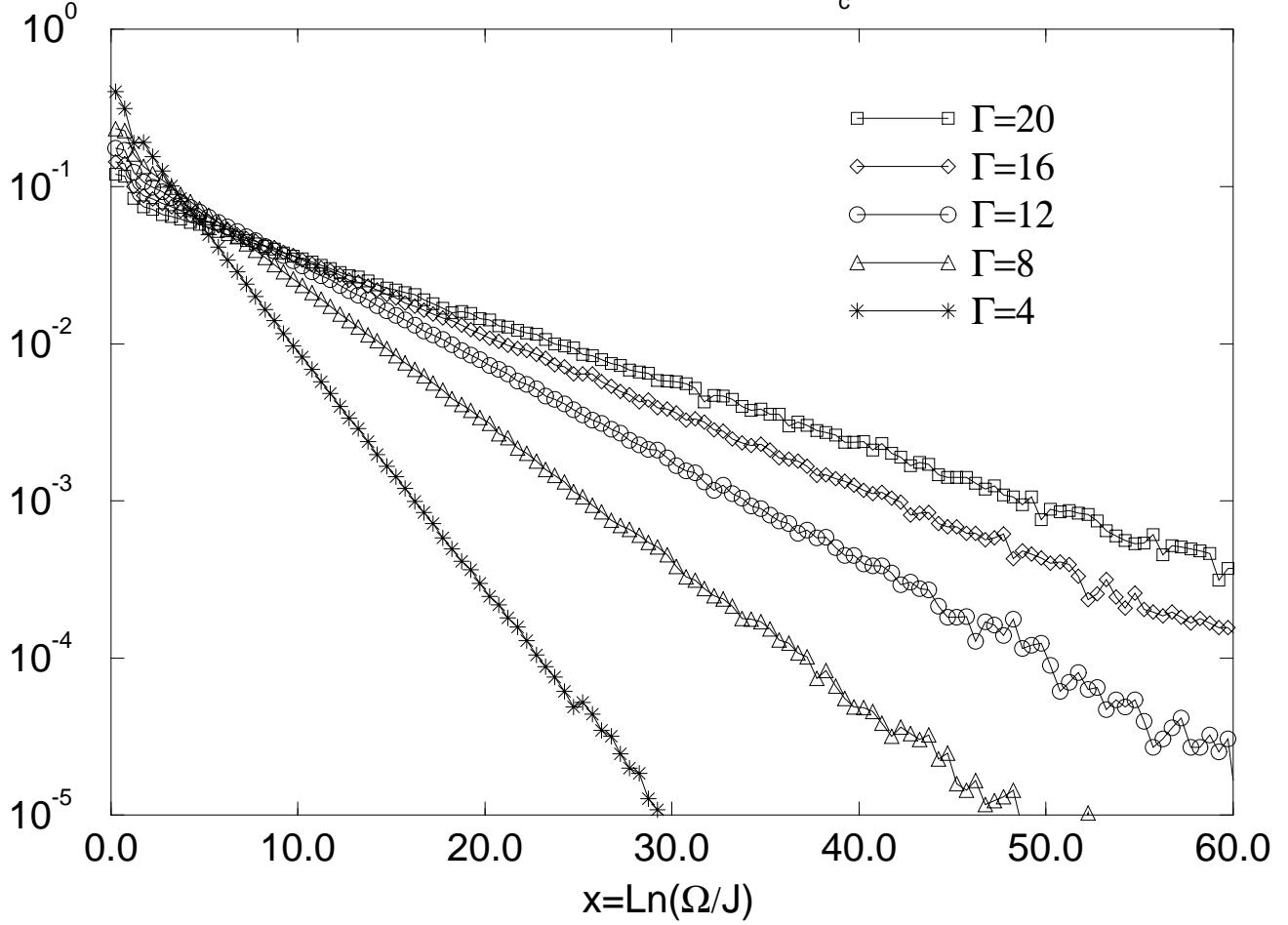


FIG. 11. Linear-Log plot of the probability distribution $P_2(x, \Gamma)$ for $\Gamma = 4, 8, 12, 16, 20$, for critical initial disorder $d_c = 5.75$: $P_2(x, \Gamma)$ takes the exponential form (5.2) with a parameter $\alpha_c(\Gamma)$ plotted on Fig 12.

$1/\alpha_c(\Gamma)$

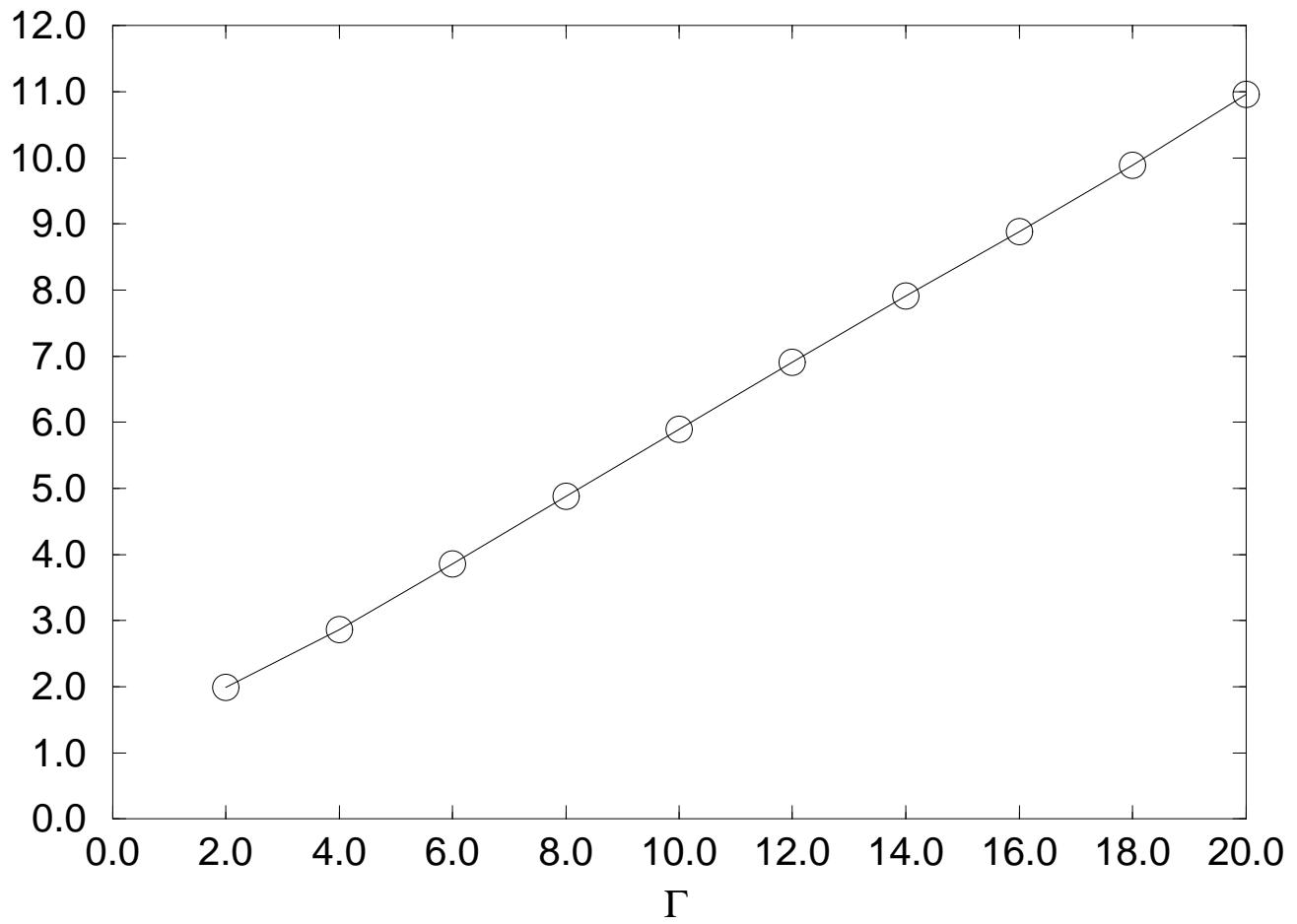


FIG. 12. Plot of the inverse of the parameter $\alpha_c(\Gamma)$ defined in (5.2) as a function of Γ : it follows the behavior (5.3).